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Semi-Exclusive Processes: New Probes of Hadron Structure

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We define and study hard "semi-exclusive" processes of the form $A + B \to C + Y$ which are characterized by a large momentum transfer between the particles A and C and a large rapidity gap between the final state particle C and the inclusive system Y. Such reactions are in effect generalizations of deep inelastic lepton scattering, providing novel currents which probe specific quark distributions of the target B at fixed momentum fraction. We give explicit expressions for photo- and leptoproduction cross sections such as $\gamma p \to \pi Y$ in terms of parton distributions in the proton and the pion distribution amplitude. Semi-exclusive processes provide opportunities to study fundamental issues in QCD, including odderon exchange and color transparency, and suggest new ways to measure spin-dependent parton distributions.

In this letter we shall study a new class of hard "semi-exclusive" processes of the form $A + B \to C + Y$, characterized by a large momentum transfer $t = (p_A - p_C)^2$ and a large rapidity gap between the final state particle C and the inclusive system Y. Here A, B and C can be hadrons or (real or virtual) photons. The cross sections for such processes factorize in terms of the distribution amplitudes of A and C and the parton distributions in the target B. Because of this factorization semi-exclusive reactions provide a novel array of generalized currents, which not only give insight into the dynamics of hard scattering QCD processes, but also allow experimental access to new combinations of the universal quark and gluon distributions.

The hard QCD processes which have been mostly studied to date can be divided into two main categories:

1. *Inclusive processes* such as DIS, $ep \to e + X$. In the limit of large photon virtuality Q^2 and energy ν in the target rest frame, the cross section can be expressed in terms of universal quark and gluon distributions $q(x, Q^2)$, $g(x, Q^2)$ in the target, where x is the fraction of target momentum carried by the struck parton and Q^2 the factorization scale.

2. *Exclusive processes* such as $ep \to ep$. For large Q^2 the form factor of the proton can be expressed in terms of its distribution amplitude, given by the valence Fock state wave function in the limit of vanishing transverse separation between the quarks [1].

More recently it has been shown that the QCD scattering amplitude for deeply virtual exclusive processes like Compton scattering $\gamma^* p \to \gamma p$ and meson production $\gamma^* p \to Mp$ factorizes into a hard subprocess and soft universal hadronic matrix elements [2,3]. For example consider exclusive meson electroproduction such as $ep \to e\pi^+ n$ (Fig. 1a). Here one takes (as in DIS) the Bjorken limit of large photon virtuality , with $x_B = Q^2/(2m_p\nu)$ fixed, while the momentum transfer $t = (p_p - p_n)^2$ re-

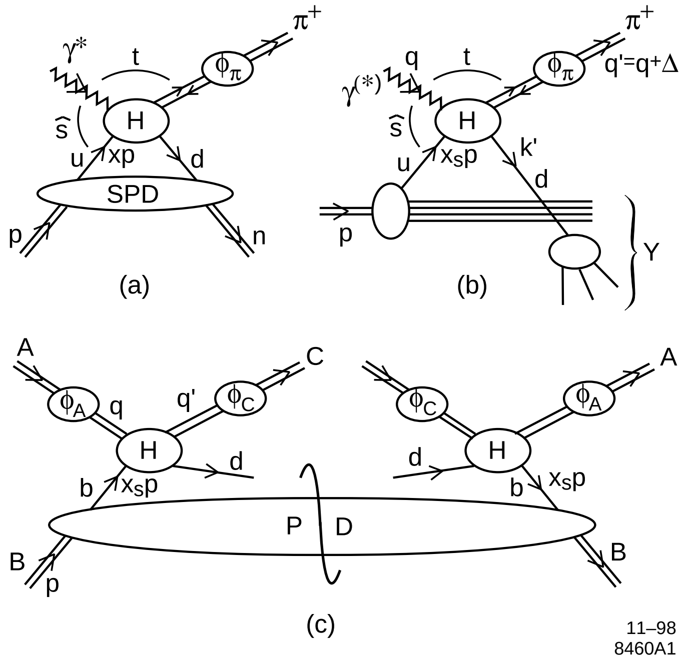

FIG. 1. (a): Factorization of $\gamma^* p \to \pi^+ n$ into a skewed parton distribution (SPD), a hard scattering H and the pion distribution amplitude ϕ_π. (b): Semi-exclusive process $\gamma^{(*)} p \to \pi^+ Y$. The d-quark produced in the hard scattering H hadronizes independently of the spectator partons in the proton. (c): Diagram for the cross section of a generic semi-exclusive process. It involves a hard scattering H, distribution amplitudes ϕ_A and ϕ_C and a parton distribution (PD) in the target B.

mains small. These processes involve 'skewed' parton distributions, which are generalizations of the usual parton distributions measured in DIS. The skewed distribution in Fig. 1a describes the emission of a u-quark from the proton target together with the formation of the final neutron from the d-quark and the proton remnants. As the subenergy $\hat{s}$ of the scattering process $\gamma^* u \to \pi^+ d$ is not fixed, the amplitude involves an integral over the u-quark momentum fraction x.

An essential condition for the factorization of the deeply virtual meson production amplitude of Fig. 1a is the existence of a large rapidity gap between the produced meson and the neutron. In fact, this factorization remains valid if the neutron is replaced with a hadronic system Y of invariant mass $M_Y^2 \ll W^2$, where W is the c.m. energy of the $\gamma^* p$ process.

For $M_Y^2 \gg m_p^2$ the momentum k' of the d-quark in Fig. 1b is large with respect to the proton remnants, and hence it forms a jet. This jet hadronizes independently of the other particles in the final state if it is not in the direction of the meson, *i.e.*, if the meson has a large transverse momentum $q'_\perp = \Delta_\perp$ with respect to the photon direction in the $\gamma^* p$ c.m. Then the cross section for an inclusive system Y can be calculated as in DIS, by treating the d-quark as a final state particle.

The large $\Delta_\perp$ furthermore allows only transversally compact configurations of the projectile A to couple to the hard subprocess H of Fig. 1b, as it does in exclusive processes [1]. Hence the above discussion applies not only to incoming virtual photons at large Q^2, but also to real photons ($Q^2 = 0$) and in fact to any hadron projectile.

Let us then consider the general process $A + B \to C + Y$, where B and C are hadrons or real photons, while the projectile A can also be a virtual photon. In the semi-exclusive kinematic limit

$$\Lambda_{QCD}^2,\, M_B^2,\, M_C^2 \ll M_Y^2,\, \Delta_\perp^2 \ll W^2 \tag{1}$$

we have a large rapidity gap

$$|y_C - y_d| = \log \frac{W^2}{\Delta_\perp^2 + M_Y^2} \tag{2}$$

between C and the parton d produced in the hard scattering (see Fig. 1c). The cross section then factorizes into the form

$$\frac{d\sigma}{dt\, dx_S}(A + B \to C + Y) = \sum_b f_{b/B}(x_S, \mu^2) \frac{d\sigma}{dt}(Ab \to Cd)\,, \tag{3}$$

where $t = (q - q')^2$ and $f_{b/B}(x_S, \mu^2)$ denotes the distribution of quarks, antiquarks and gluons b in the target B. The momentum fraction x_S of the struck parton b is fixed by kinematics to the value

$$x_S = \frac{-t}{M_Y^2 - t} \,, \tag{4}$$

and the factorization scale μ^2 is characteristic of the hard subprocess $Ab \to Cd$.

In the kinematic limit (1) the subenergy $\hat{s} = (q + x_S p)^2$ of the hard process, the momentum transfer t, and the fraction x_F of the light-cone momentum of projectile A carried by particle C are respectively given by

$$\hat{s} = \frac{\Delta_\perp^2}{\Delta_\perp^2 + M_Y^2} W^2 \,, \tag{5}$$

$$-t = \frac{\Delta_\perp^2 + x_B M_Y^2}{1 - x_B} = \frac{\Delta_\perp^2}{1 - x_B/x_S} \,, \tag{6}$$

$$1 - x_F = \frac{\Delta_\perp^2 + M_Y^2}{W^2} \,, \tag{7}$$

where we notice that x_F is close to 1. We also have the relation

$$x_B < x_S = \frac{\Delta_\perp^2 + x_B M_Y^2}{\Delta_\perp^2 + M_Y^2} < 1 \,, \tag{8}$$

with $x_B = 0$ in the case where the projectile A is a hadron or real photon.

It is conceptually helpful to regard the hard scattering amplitude H in Fig. 1c as a generalized current of momentum $q - q' = p_A - p_C$, which interacts with the target parton b. For $A = \gamma^*$ we obtain a close analogy to standard DIS when particle C is removed. With $q' \to 0$ we thus find $-t \to Q^2$, $M_Y^2 \to W^2$, and see that x_S in (4) goes over into $x_B = Q^2/(W^2 + Q^2)$. The possibility to control the value of q' (and hence the momentum fraction x_S of the struck parton) as well as the quantum numbers of particles A and C should make semi-exclusive processes a versatile tool for studying hadron structure. The cross section further depends on the distribution amplitudes ϕ_A, ϕ_C (*cf.* Fig. 1c), allowing new ways of measuring these quantities. The use of this new current requires a sufficiently high c.m. energy, since according to Eq. (1) we need to have at least one intermediate large scale. We note that the possibility of creating effective currents using processes similar to the ones we discuss here was considered already before the advent of QCD [4].

It is instructive to compare our semi-exclusive limit (1) for electroproduction, $Q^2 \sim W^2$, with the $x_F \to 1$ limit of semi-inclusive DIS. After being struck by the virtual photon the u-quark in Fig. 1b has a virtuality $\hat{s} \sim Q^2$ when $\Delta_\perp^2 \sim M_Y^2$, cf. Eq. (5). Since the time scale $1/Q$ of the photon interaction is then similar to the time scale $1/\sqrt{\hat{s}}$ of the further interactions of the struck quark, these processes cannot be physically separated. Hence the hard subprocess H of Fig. 1b is compact. On the other hand, if $\Delta_\perp^2 \ll M_Y^2$ the virtual photon time scale is much shorter than that of quark fragmentation, and H factorizes into $\gamma^* u \to u$ times $u \to \pi^+ d$. This is the physics of semi-inclusive DIS and also of lepton pair production ($\pi p \to \mu^+ \mu^- Y$) in the limit $x_F \to 1$ when $\Delta_\perp$ is integrated over [5].

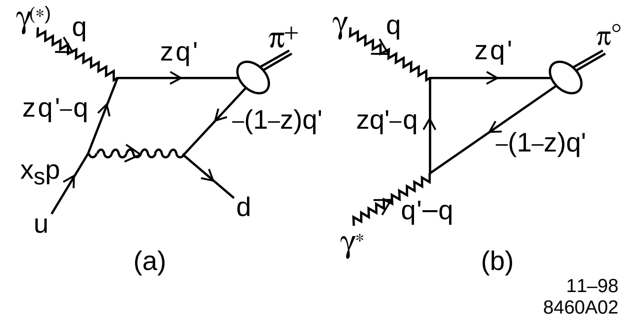

FIG. 2. (a): One of the diagrams of the hard scattering $\gamma^{(*)} u \to \pi^+ d$ to leading order in α_s. (b): One of the leading order diagrams for the pion transition form factor in $\gamma^* \gamma \to \pi$.

Pion photoproduction at large transverse momentum $\Delta_\perp$ was studied in Ref. [6] for $x_F < 1$, *i.e.*, in the case of no rapidity gap. In this case the struck quark emits a gluon at a short time-scale $1/\Delta_\perp$, but the pion is predominantly produced via a standard non-perturbative fragmentation process.

Next we consider in more detail the specific semi-exclusive process $\gamma^{(*)} p \to \pi^+ Y$ shown in Fig. 1b. We work in the kinematic limit (1) and for simplicity take a single intermediate scale, $\Delta_\perp^2 \sim M_Y^2$. The virtuality Q^2 of the photon can scale as

$$Q^2 \sim \begin{cases} 0 & \text{photoproduction} \\ M_Y^2 & \text{DIS, } x_B \to 0 \\ W^2 & \text{DIS, } x_B \text{ finite} \end{cases} \tag{9}$$

Note that according to Eqs. (5) and (6) $-t \sim \Delta_\perp^2$ is of intermediate scale, and $\hat{s} \sim W^2$ is very large, so that we have $\Lambda_{QCD}^2 \ll -t \ll \hat{s}$ in the hard scattering.

The target parton (b in Fig. 1c) attached to the hard amplitude H can (at lowest order in α_s) be either one of the valence quarks $u, \bar{d}$ of the π^+. These two contributions add incoherently in the cross section, weighted by the respective parton distribution $f_{b/p}$. In Fig. 2a we show one the four diagrams which contribute to H in the case $b = u$. The three other diagrams are obtained by different orderings of the photon and gluon vertices on the u- and d-quark lines.

Exclusive processes can be sensitive to infrared end-point contributions, where the momentum of one of the valence quarks in a hadron wave function vanishes. In Fig. 2a it may be seen that the gluon propagator goes on-shell for $z \to 1$, since its momentum then equals that of the final d-quark. It is important to note that the gluon four-momentum does not vanish in this limit, since it still has a large transverse component $-\vec{\Delta}_\perp$. As a consequence the u-quark propagator does not become singular at this end-point, in other words $(zq' - q)^2 = zt - (1-z)Q^2$ is "protected" by the large momentum transfer $-t$ in the

$z \to 1$ limit. This suggests that the hard amplitude is no more sensitive to end-point contributions than the pion transition form factor (Fig. 2b). On the other hand, in exclusive meson production at large Q^2 and small t both internal propagators in Fig. 2a go on-shell at $z = 1$, which is what makes the amplitude with transversely polarized virtual photons infrared sensitive [3].

We also note that (for $z \neq 0, 1$) all propagators in the hard scattering subprocess have at least a virtuality of order Q^2 or $-t$, whichever is larger. This ensures that the scattering amplitude H is compact, and that the photon couples coherently to both valence quarks (*i.e.*, all four diagrams contribute at leading order). In contrast to ordinary DIS and semi-inclusive processes, the contribution of the parton distribution $f_{u/p}$ will thus not necessarily be weighted by e_u^2, the square of the electric charge of the struck quark.

In the case of photoproduction one finds in the limit of Eq. (1) with $Q^2 = 0$ that the $\gamma u \to \pi^+ d$ subprocess cross section is [6]

$$\frac{d\sigma}{dt}(\gamma u \to \pi^+ d) = \frac{128\pi^2}{27}\, \alpha\alpha_s^2\, \frac{(e_u - e_d)^2}{\hat{s}^2(-t)} \times \left\{ \left[\int_0^1 dz \frac{\phi_\pi(z)}{z}\right]^2 + \left[\int_0^1 dz \frac{\phi_\pi(z)}{1-z}\right]^2 \right\}\,, \quad (10)$$

where we choose the convention of Refs. [1,6] for the pion distribution amplitude, namely $\int dz\, \phi_\pi(z) = f_\pi/\sqrt{12}$ with $f_\pi = 93$ MeV. The $\gamma \bar{d} \to \pi^+ \bar{u}$ subprocess has an identical cross section. The result for the physical process $\gamma p \to \pi^+ Y$ is then, according to Eq. (3),

$$\frac{d\sigma}{dt\, dx_S}(\gamma p \to \pi^+ Y) = \left[u(x_S, -t) + \bar{d}(x_S, -t)\right] \frac{d\sigma}{dt}(\gamma u \to \pi^+ d) \qquad (11)$$

with the notation $q(x_S, -t) = f_{q/p}(x_S, -t)$. There are several interesting aspects of this result.

- Both the target u- and $\bar{d}$-quark contributions are weighted by the *total charge* $e_u - e_d = +1$ of the produced π^+. An analogous formula holds of course if the π^+ is replaced with another pseudoscalar. For neutral meson production, $\gamma p \to M^0 Y$ with $M^0 = \pi^0, K^0, \eta, \ldots$, the expression (10) *vanishes*; more exactly one finds that the cross section is suppressed by $(-t/\hat{s})^2$ compared with the charged meson case [6]. This illustrates that the new type of current probe we are considering weights the target parton distributions differently from the photon current, as is also the case for weak currents.

- The cross section Eq. (10) has a power-law behavior, $d\sigma/dt \propto 1/\hat{s}^3$ at fixed $t/\hat{s}$. This is the basic signature that the amplitude factorizes into a meson distribution amplitude and a hard scattering subprocess.

- At fixed t the expression (10) goes like $1/\hat{s}^2$, which is characteristic of two spin 1/2 quark exchanges in the t-channel. Notice that with $\Lambda_{QCD}^2 \ll -t \ll \hat{s}$ the hard scattering takes place in the perturbative Regge regime: if a deviation from this $\hat{s}$-behavior were to be observed experimentally it would indicate that the quark exchange reggeizes, *i.e.*, that contributions from higher order ladder diagrams are important. For a discussion of QCD expectations and experimental evidence on how meson Regge trajectories $\alpha(t)$ behave at large $-t$ we refer to [7]. Let us emphasize that perturbative Reggeization would not change the power-law behavior in $\hat{s}$ at fixed $t/\hat{s}$ mentioned above.

- The pion distribution amplitude $\phi_\pi(z)$ enters in precisely the same way as it does in the pion transition form factor for $\gamma^*\gamma \to \pi^0$ [1,6]

$$F_{\pi\gamma}(Q^2) = \frac{\sqrt{48}\,(e_u^2 - e_d^2)}{Q^2} \int_0^1 dz \frac{\phi_\pi(z)}{z}\,. \qquad (12)$$

  A comparison of Figs. 2a and 2b suggests that this may again be interpreted as the result of replacing the γ^* probe with an effective $u\bar{d}$ current of hardness t. Hence the relation between observables,

$$\frac{d\sigma}{dt}(\gamma u \to \pi^+ d) = \frac{16\pi^2}{9}\, \alpha\alpha_s^2\, \frac{-t|F_{\pi\gamma}(-t)|^2}{\hat{s}^2}\,, \quad (13)$$

  which holds at lowest order in α_s, may have a broader range of validity. Note that in order to obtain Eq. (13) we have used isospin symmetry, $\phi_{\pi^+}(z) = \phi_{\pi^0}(z)$, which also implies $\phi_\pi(z) = \phi_\pi(1-z)$.

In the limit (1) with $Q^2 \sim W^2$, *i.e.*, at finite x_B, the semi-exclusive electroproduction cross section is

$$\frac{d\sigma(ep \to e\pi^+ Y)}{dQ^2\, dx_B\, dt\, dx_S} = \frac{\alpha}{\pi} \frac{1-y}{Q^2 x_B} \frac{512\pi^2}{27}\, \alpha\alpha_s^2 \frac{x_B}{\hat{s}\, Q^4 x_S} \times \left[\int_0^1 dz \frac{\phi_\pi(z)}{z}\right]^2 \left\{ u(x_S) \left[e_u + \left(1 - \frac{x_B}{x_S}\right) e_d\right]^2 + \bar{d}(x_S) \left[e_d + \left(1 - \frac{x_B}{x_S}\right) e_u\right]^2 \right\} \quad (14)$$

where $y = \nu/E_e$ is the momentum fraction of the projectile electron carried by the virtual photon, and we have used again $\phi_\pi(z) = \phi_\pi(1-z)$. We make the following remarks.

- The semi-exclusive cross section in Eq. (14) corresponds to longitudinal photon exchange. The contribution from transverse photons is suppressed, as in the exclusive case $\gamma^* p \to Mp$ at large Q^2 and small $-t$ [3].

- For $\Delta_\perp^2 \ll M_Y^2$ we have $x_S \to x_B$ according to Eq. (8). We find that the parton distributions are then multiplied by the corresponding quark charge squared in the cross section. This is a consequence of the fact that, as discussed above, the hard subprocess factorizes in this limit into a virtual photon interaction and a quark fragmentation process. We note that Eq. (14) was derived for $\Delta_\perp^2 \sim M_Y^2$ and acquires corrections when $\Delta_\perp^2 \ll M_Y^2$.

- The subprocess is of hardness Q^2, which thus is the relevant scale for the quark parton distributions. If one takes into account higher orders in α_s the target parton couples to a ladder. Its hardness, which then becomes the appropriate scale, will be between $-t$ and Q^2.

In the intermediate range $Q^2 \sim M_Y^2$ of Eq. (9) transverse and longitudinal photon polarizations contribute with comparable strength, and the structure of the cross section is richer than in the two extreme cases just discussed.

We conclude with a number of more general remarks and suggestions for future work.

1. Vector mesons. In addition to pseudoscalar mesons one can also consider vector meson production. We find that exact analogs of Eqs. (10) and (14) hold if the vector meson is longitudinally polarized. Transverse vector mesons are suppressed in the cross section by $(-t/\hat{s})^2$ for photo- and $-t/\hat{s}$ for electroproduction. For symmetry reasons there is no interference between different meson polarizations.

2. Particle production ratios. Systematic comparisons of semi-exclusive photoproduction of various particles can give useful information on parton distributions and distribution amplitudes. The hard subprocess (10) cancels in the ratio of physical cross sections (11) for π^+ and π^-. Hence $d\sigma(\pi^+)/d\sigma(\pi^-)$ directly measures the $(u+\bar{d})/(d+\bar{u})$ parton distribution ratio. Similarly, the $d\sigma(K^+)/d\sigma(K^-)$ ratio measures the strange quark content of the target without uncertainties due, *e.g.*, to fragmentation functions.

Conversely, the parton distributions drop out in the ratio $d\sigma(\rho_L^+)/d\sigma(\pi^+)$ of longitudinally polarized ρ mesons to pions, allowing a comparison of their distribution amplitudes. Since the normalization of both ϕ_ρ and ϕ_π is fixed by the leptonic decay widths such a comparison can reveal differences in their z-dependence. In the intermediate Q^2-range of Eq. (9) the relative size of Q^2 and t can furthermore be tuned to change the dependence of the hard subprocess on z. This may be used to get further information on the shape of $\phi(z)$.

3. Hard odderon and pomeron exchange. As we noted above, the lowest order quark exchange contribution to $\gamma p \to \pi^0 Y$ is strongly suppressed. At higher orders in α_s there is a contribution from scattering on gluons ($b = g$ in Fig. 1c). Due to charge conjugation at least two gluons need to be emitted from the hard scattering ($d = gg$ in Fig. 1c). Altogether three gluons are thus exchanged in the t-channel, corresponding to hard odderon exchange. Semi-exclusive production of π^0, η or other neutral pseudoscalars may thus be particularly sensitive to odderon effects since competing production mechanisms are suppressed. See Ref. [8] for a recent discussion of odderon physics.

An analogous argument suggests that $\gamma p \to \rho^0 Y$ is a good process for studying hard two-gluon ladders, *i.e.*, pomeron exchange at large t. In this context neutral vector meson production has in fact been studied in a kinematic region very similar to ours [9,10]. Real photon production, $\gamma p \to \gamma Y$, may also be interesting in this respect [11], since compared to two-gluon exchange the quark exchange contribution is again suppressed by a power of $-t/\hat{s}$.

4. Spin and transversity distributions. Polarization of the target B can naturally be incorporated in our framework. A longitudinally polarized target selects the usual spin-dependent parton distributions $\Delta q(x_S)$. It also appears possible to measure the quark transverse spin, or transversity distribution in photoproduction of ρ mesons on transversely polarized protons. In this case only the interference term between transversely and longitudinally polarized ρ mesons should contribute. Since very little experimental information on the transversity distribution is available, this question merits further study.

5. Color transparency. The factorization of the hard amplitude H in Fig. 1c from the target remnants is a consequence of the high transverse momentum which selects compact sizes in the projectile A and in the produced particle C. In the case of nuclear targets this color transparency [12] implies according to Eq. (3) that all target dependence enters via the nuclear parton distribution. Thus tests of color transparency can be made even in photoproduction, *e.g.*, through

$$\gamma A \to \left\{ \begin{array}{c} \pi^+(\Delta_\perp) + Y \\ p(\Delta_\perp) + Y \end{array} \right. \tag{15}$$

in the semi-exclusive kinematic region specified by Eq. (1).

Color transparency has so far been studied mainly in exclusive processes where the target scatters elastically, such as $\gamma^* A \to \rho A$ [13], $pA \to pp + (A-1)$ [14] and $\gamma^* A \to p + (A-1)$ [15]. Semi-exclusive processes provide a possibility to study color transparency in processes where the target dissociates into a heavy inclusive system Y [16]. This puts less stringent requirements on the energy resolution of the apparatus, but it also requires a higher beam energy to ensure the existence of a rapidity gap.

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